HIGGS SEARCH AT e^+e^- AND $\gamma\gamma$ COLLIDERS[*]

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ABSTRACT

The prospects for discovering Higgs particles and studying their fundamental properties at future high–energy electron–positron and photon-photon colliders are reviewed. Both the Standard Model Higgs boson and the Higgs particles of its minimal supersymmetric extension are discussed. We also comment on a two Higgs doublet model searches at LEP 1. We update various results by taking into account the value of the top quark mass obtained by the CDF Collaboration and by including radiative corrections.

1. Introduction

The Higgs mechanism of electroweak symmetry breaking is a crucial ingredient of the Standard Model ($\mathcal{SM}$). It allows to give masses to weak gauge bosons without affecting the renormalizability of the model. At the same time fermions get their masses via Yukawa interaction with the ground state of the Higgs field. This goal is achieved by employing one iso-doublet of scalar Higgs fields which leads to one physical Higgs boson.[1] The fundamental nature of mass generation of quarks, leptons and weak gauge bosons requires this picture to be verified experimentally in all its aspects. Therefore the Higgs search constitues one of the most important physics goals at current and future experiments and is one of the motivations for building new accelerators.[2,3,4,5] If the Higgs particle is found its properties must be measured, in particular the couplings to other particles which are uniquely fixed by Higgs mechanism.

Although one can argue that it is a matter of time until the Higgs boson of the $\mathcal{SM}$ is found, some theoretical problems related to its existence suggest that it is necessary to look beyond the Standard Model. Scalar fields have a nice property that thay can have a nonzero vacuum expectation value without breaking Lotentz invariance and thus can trigger breakdown of gauge symmetry. However they also have a bad property of aquiring quadratic divergences through radiative corrections. The correction to the mass of the Higgs boson is $\delta m^2 \sim g^2 \Lambda^2$, where Λ is a physical scale beyond which the low energy theory no longer applies. To understand the

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physics at the Fermi scale $\sim 10^2$ GeV it would be inappropriate to have Λ of the order of the unification scale $\sim 10^{15}$ GeV or the Planck scale $\sim 10^{19}$ GeV but rather in the TeV region. This is so called the naturalness or hierarchy problem. Therefore we are led to consider means to stabilize the Higgs sector below or at TeV scale. The only way to protect masses of elementary scalar particles is supersymmetry.[6] This symmetry relates bosons to fermions and makes the bosons to behave as well as fermions.

In this review we start in section 2 with the presentation of the Higgs sector in $\mathcal{SM}$. Various decay modes of the Higgs boson and its production mechanisms at current and future e^+e^- and $\gamma\gamma$ colliders with an energy in the range 300 - 500 GeV are discussed. In section 3 we discuss the minimal supersymmetric ($\mathcal{MSSM}$) extension of the model in which two Higgs doublets are present. Although some features are specific to this version of the model, the pattern of phenomena to be discussed is characteristic to more general SUSY models. We will also comment on Higgs search in a two Higgs doublet model ($2\mathcal{HDM}$) at LEP 1. Section 4 contains conclusions and outlook.

We discuss Higgs searches in e^+e^- and $\gamma\gamma$ collisions. The Higgs search at hadron colliders has been discussed in the talk of Poggioli.[7] We update various results for Higgs masses and couplings, partial decay widths and production cross sections taking into account the value of the top quark mass $m_t = 175$ GeV consistent with the value recently published by the CDF collaboration[8] and with the favored value obtained from a global fit[9] of electroweak precision measurements at LEP and SLC. Sometimes we will take the values 150 or 200 GeV which can be viewed as conservative lower and upper bounds on the top mass, respectively.

2. Search for the Higgs boson of the Standard Model

2.1. The Higgs sector of the $\mathcal{SM}$

In the Standard Model the mass of the Higgs boson is the only unknown parameter. Once the mass is fixed, the profile of the Higgs particle can be predicted completely.

Even though the value of the Higgs mass cannot be predicted, interesting constraints can be derived from assumptions: a) on the energy range within which the model is valid before perturbation theory breaks down at a scale Λ and new dynamical phenomena would emerge, and b) stability of the vacuum.
a) Since the strength of the Higgs self–interaction is determined by the Higgs mass itself M_H, the condition $M_H < \Lambda$ sets an upper limit on the Higgs mass in the Standard Model. Thorough analyses lead to an estimate of about 630 GeV for the upper limit of M_H. On the other hand, if the Higgs mass is less than 180 to 200 GeV, the Standard Model can be extended up to the GUT scale $\Lambda_{\mathrm{GUT}} \sim 10^{15}$ GeV with weakly interacting particles. Including the effect of t–quark loops on the running coupling, a detailed analysis predicts the area of the allowed (m_t, M_H) values shown in Fig. 1 for several values of the cut–off parameter Λ.[10]
b) Quantum corrections due to top-quark loops to the quartic Higgs coupling are

negative, driving the coupling to negative values for which the vacuum becomes unstable. For top masses larger than about 100 GeV, this leads to a lower limit on the Higgs mass, Fig. 1.

From the above arguments the Higgs mass could well be expected in the window $100 < M_H < 180$ GeV for a top mass value of 150 GeV and $160 < M_H < 200$ GeV for $m_t \sim 175$ GeV. The mass range $M_Z - 2M_Z$ is usually referred to as intermediate Higgs mass range.

Fig. 1 Values of the top and Higgs masses for which the Standard Model can be extended up to the scale Λ; from Ref.[10] The lower bound is derived from vacuum stability. The dashed line shows the upper limit for the mass og the lighest scalar Higgs boson of $\mathcal{MSSM}$, see Section 3 and Fig. 6

2.2. Decay modes

The strength of the couplings of the Higgs boson to fermions and to the electroweak gauge bosons $V = W, Z$ is set by their masses:

$$g_{ffH} = \left[\sqrt{2}G_F\right]^{1/2} m_f \tag{1}$$

$$g_{VVH} = 2\left[\sqrt{2}G_F\right]^{1/2} M_V \tag{2}$$

In the Born approximation the width of the Higgs decay into fermion pairs is

$$\Gamma(H \to f\bar{f}) = N_c \frac{G_F m_f^2}{4\sqrt{2}\pi} M_H \, \beta^3 \tag{3}$$

with $\beta = (1 - 4m_f^2/M_H^2)^{1/2}$, $N_c = 1$ for leptons and $N_c = 3$ for quarks. For quark pairs, one has to use running quark masses evaluated at the scale $\mu = M_H$[11] which

take the bulk of QCD corrections. For example, in the case of the b quark and for Higgs masses around 100 GeV, it amounts to a reduction of the $H \to b\bar{b}$ decay width by more than 50%.

Above the $H \to WW$ and ZZ decay thresholds, the partial width into massive gauge boson pairs may be written as[12]

$$\Gamma(H \to VV) = \delta_V \frac{\sqrt{2}G_F}{32\pi} M_H^3 (1 - 4x + 12x^2)\beta \,, \tag{4}$$

where $x = M_V^2/M_H^2$, $\beta = \sqrt{1 - 4x}$ and $\delta_V = 2(1)$ for $V = W(Z)$.

Below the threshold for two real bosons, the Higgs particle can decay into real and virtual VV^* pairs (primarily WW^* pairs above $M_H \sim 110$ GeV) followed by V^* decay into a fermion pair. The partial decay width is given by[13]

$$\Gamma(H \to VV^*) = \frac{3G_F^2 M_V^4}{16\pi^3} M_H R(x) \delta_V' \,, \tag{5}$$

with $\delta_W' = 1$ and $\delta_Z' = 7/12 - 10\sin^2\theta_W/9 + 40\sin^4\theta_W/27$ and

$$R(x) = \frac{3 - 24x + 60x^2}{(4x - 1)^{1/2}} \arccos\frac{3x - 1}{2x^{3/2}} - \frac{1 - x}{2x}(2 - 13x + 47x^2) - (\frac{3}{2} - 9x + 6x^2)\log x$$

Fig. 2 Total decay width (a) and decay branching ratios (b) of the $\mathcal{SM}$ Higgs boson; the top quark mass is fixed to $m_t = 175$ GeV. The QCD corrections to the hadronic decay modes are included.

By adding up all possible decay channels we obtain the total width shown in Fig. 2a for $m_t = 175$ GeV. The Higgs particle is very narrow ($\Gamma(H) \leq 10$ MeV) below the (virtual) gauge boson channels and becomes rapidly wider reaching ~ 1 GeV at the ZZ threshold. Only above $M_H \geq 250$ GeV it becomes wide enough to be resolved experimentally.

The branching ratios of the main decay modes are displayed in Fig. 2b from Ref.[14], an update of Ref.[15]. A large variety of channels will be accessible for Higgs masses below 140 GeV. By far the dominant mode are $b\bar{b}$ decays, yet $c\bar{c}$, $\tau^+\tau^-$ and gg still occur at a level of several percent. The branching ratios for the $H \to \gamma\gamma$ and γZ are small, being of $\mathcal{O}(10^{-3})$. Above the mass value $M_H = 140$ GeV, the Higgs boson decay into W's becomes dominant, overwhelming all other channels once the decay mode into two real W's is kinematically possible.

Although the coupling $H\gamma\gamma$ is small, it is very interesting since it is sensitive to all charged particles in the loop and can be used as a possible probe of new particles whose masses are generated by the Higgs mechanism. This will be important when we will discuss Higgs production in $\gamma\gamma$ collisions.

2.3. SM Higgs search at LEP

Prior to LEP only the Higgs masses up to ~ 5 GeV have been excluded. The most comprehensive search for Higgs particles has been carried out in Z decays at LEP,[16] based on the Bjorken process $Z \to Z^*H$. Both visible $Z^* \to q\bar{q}$ and l^+l^- and invisible $Z^* \to \nu\bar{\nu}$ decay modes of the virtual Z in processes with Higgs boson decays into $b\bar{b}$, $c\bar{c}$ and $\tau^+\tau^-$ have been looked for. By now, a lower bound on the Higgs mass of about $M_H \geq 63.5$ GeV can be established. This limit can be raised by a few GeV by accumulating more statistics. At $M_H \sim 70$ GeV another production process $Z \to H\gamma$ can also be explored. In the second phase of LEP with a total energy close to 200 GeV Higgs particles can be searched for up to masses of ~ 110 GeV in Higgs bremsstrahlung off the Z line. In the range $m_Z \pm 10$ GeV the sensitivity will be weaker due to irreducible background coming from $e^+e^- \to ZZ$ and the b quark tagging will be necessary to discover or exclude the Higgs. Higher energy colliders are required to sweep the entire mass range for the Higgs particle.

2.4. Production mechanisms at future e^+e^- colliders

At e^+e^- linear colliders operating in the 300–500 GeV energy range, the main production mechanisms for Higgs particles are the following processes:[4]

$$\text{bremsstrahlung process}: \quad e^+e^- \longrightarrow (Z) \longrightarrow Z + H \tag{6}$$

$$\text{WW fusion process}: \quad e^+e^- \longrightarrow \bar{\nu}\,\nu\,(WW) \longrightarrow \bar{\nu}\,\nu\, + H \tag{7}$$

$$\text{ZZ fusion process}: \quad e^+e^- \longrightarrow e^+e^-(ZZ) \longrightarrow e^+e^- + H \tag{8}$$

$$\text{radiation off top}: \quad e^+e^- \longrightarrow (Z,\gamma) \longrightarrow (t\bar{t}) \longrightarrow t\bar{t} + H \tag{9}$$

The mass dependence of the cross sections at $\sqrt{s} = 500$ GeV is shown in Fig. 3.

For M_H in the range $150 - 200$ GeV the cross sections for the bremsstrahlung and the WW fusion processes are of comparable size at $\sqrt{s} = 500$ GeV, while the

ZZ fusion cross section is smaller by an order of magnitude. With $\sigma \sim 100$ fb, a total of ~ 2000 Higgs particles per year can be created at an integrated luminosity of $\int \mathcal{L} = 20$ fb^{-1}. For Higgs masses below 100 GeV, the cross section for Higgs radiation off top quarks Eq. (9) is of the order of a few fb;[17] this process can be used only to measure the $t\bar{t}H$ Yukawa coupling once the Higgs boson is detected in the previous processes.

Fig. 3 Production cross sections for $\mathcal{SM}$ Higgs particles at $\sqrt{s} = 500$ GeV.

2.4.1. Higgs Bremsstrahlung

The cross section for the process Eq. (6) can be written in a compact form

$$\sigma(e^+e^- \to ZH) = \frac{G_F^2 M_Z^4}{96\pi s}(v_e^2 + a_e^2)\ \lambda^{1/2}\frac{\lambda + 12M_Z^2/s}{(1 - M_Z^2/s)^2} \tag{10}$$

where $a_e = -1$ and $v_e = -1 + 4s_W^2$ are the Z charges of the electron and $\lambda = (1 - M_H^2/s - M_Z^2/s)^2 - 4M_H^2 M_Z^2/s^2$ is the usual two–particle phase space function.

Due to kinematical constraints familiar from low-energy e^+e^- experiments there are several strategies that can be adopted. First, the recoiling Z boson in the two–body reaction $e^+e^- \to ZH$ is mono–energetic and the mass of the Higgs boson can be derived from the energy of the Z boson, $M_H^2 = s - 2\sqrt{s}E_Z + M_Z^2$. Second, the Higgs boson can be reconstructed from its decay products with an additional constraint provided by the recoiling system which should have an invariant mass to be equal M_Z. Finally full reconstruction of the events can be attempted.

In the first two cases the initial e^+ and e^- beam energies have to be well known. However, beamstrahlung smears out the c.m. energy and the system moves along the beam axes. The intensity of the beamstrahlung depends on the machine design. In this context promising results[18] have been obtained in the DESY–Darmstadt and the TESLA design studies. For these designs the smearing of the missing mass is expected to be of the same magnitude as the experimental uncertainties in the reconstruction of the Z boson in the leptonic decay channels.

Since the recoiling Z boson remains approximately mono–energetic, even if beam-strahlung is taken into account, it is easy to separate the signal from the background[19]. Only for Higgs masses close to Z mass the selection of $b\bar{b}$ final states from H decays by means of flavor tagging through vertex detection will be necessary in order to eliminate the dominant background from double Z–production $e^+e^- \to ZZ$. For example, at $\sqrt{s} = 500$ GeV and integrated luminosity $\int \mathcal{L} = 20$ fb^{-1} one expects 50 ZH and 142 ZZ events without b-tagging, and 38 and 33 with b-tagging, respectively. For masses M_H between 100 and 160 GeV the dominant background from single Z–production in $e^+e^- \to ZZ^*(\to q\bar{q})$ and $e^+e^- \to Z + WW^*(\to q\bar{q}')$ is suppressed by at least one power of the electroweak coupling relative to the signal. Beyond 160 GeV and 180 GeV, the reactions with three gauge bosons in the final state, $e^+e^- \to Z + WW$ and $e^+e^- \to Z + ZZ$ are the main background channel with the invariant mass of the WW or ZZ final states broad as opposed to the resonance structure of the signal and missing mass technique can be used up to kinematical limit.

An interesting feature of the bremsstrahlung process is that the angular distribution of the Z/H bosons is sensitive to the spin of the Higgs particle.[20] At high energies the Z boson in Eq. (6) is produced with longitudinal polarization and the angular distribution approaches the spin–zero angular distribution asymptotically $\mathrm{d}\sigma/\mathrm{d}\cos\theta \to \frac{3}{4}\sin^2\theta$. For a pseudoscalar state $A(0^{-+})$ (realized for example in $2\mathcal{HDM}$ and $S\tilde{U}SY$ models) the effective point-like coupling ZZA is a P–wave coupling, as opposed to a S–wave for a scalar Higgs, and in this case the angular distribution is $\mathrm{d}\sigma(ZA)/\mathrm{d}\cos\theta \sim 1 - \frac{1}{2}\sin^2\theta$, independent of the energy. The angular distributions specific to Higgs production in $e^+e^- \to ZH$ or $e^+e^- \to ZA$ are different from the process $e^+e^- \to ZZ$ which is mediated by electron exchange in the t–channel and the amplitude is built up by many partial waves, peaking in the forward/backward direction. The three distributions (ZH, ZA, ZZ) are compared with each other in Fig. 4, which demonstrates the specific character of the Higgs production process.

Since the Z bosons from ZH production are expected to be asymptotically polarized longitudinally (by contrast, the Z bosons from ZA associate production or ZZ pair production are transversally polarized) this pattern can be checked further experimentally by studying the distribution of the light fermions in the $Z \to f\bar{f}$ rest frame.

Fig. 4 Angular distributions of the processes $e^+e^- \to ZH, ZA$ and ZZ at $\sqrt{s} = 500$ GeV. The Higgs mass is fixed to $M_H = 120$ GeV.

2.4.2. WW and ZZ Fusion Processes

The cross section for the fusion processes, Eq. (7,8) can be written as

$$\sigma = \frac{G_F^3 M_V^4}{64\sqrt{2}\pi^3} \int_{\kappa_H}^1 \mathrm{d}x \int_x^1 \frac{\mathrm{d}y}{[1+(y-x)/\kappa_V]^2} \left[(v^2+a^2)^2 f(x,y) + 4v^2a^2 g(x,y)\right] \quad (11)$$

$$
\begin{aligned}
f(x,y) &= \left(\frac{2x}{y^3} - \frac{1+2x}{y^2} + \frac{2+x}{2y} - \frac{1}{2}\right)\left[\frac{z}{1+z} - \log(1+z)\right] + \frac{x}{y^3}\frac{z^2(1-y)}{1+z} \\
g(x,y) &= \left(-\frac{x}{y^2} + \frac{2+x}{2y} - \frac{1}{2}\right)\left[\frac{z}{1+z} - \log(1+z)\right]
\end{aligned}
$$

with $\kappa_H = M_H^2/s, \kappa_V = M_V^2/s, z = y(x-\kappa_H)/(\kappa_V x)$ and v, a the electron couplings to the massive gauge bosons $[v = -1+4s_W^2, a = -1$ for the Z boson and $v = a = \sqrt{2}$ for the W boson].

Asymptotically it grows as $M_W^{-2}\log s/M_H^2$ in contrast to the processes with s-chaneel exchanges which fall $\sim s^{-1}$. The cross section for ZZ fusion is about an order of magnitude smaller than the cross section for WW fusion; this is a mere consequence of the fact that the NC couplings are smaller than the CC couplings. The lower rate however is, at least partly, compensated by the clean signature of the e^+e^- final state in (8) that allows for a missing mass analysis to tag the Higgs particle.

For a light Higgs mass, the dominant background comes from $e^+e^- \to e^+W^-\nu_e$ and $WW \to Z$ processes. The cross sections exceed the signal for jet–jet final states by about a factor of 60 and 3, respectively. Since the single $W(Z)$ boson production shows a behavior similar to the Higgs boson in the signal process, kinematical cuts enhance the signal/background ratio very little. Only the use of features that are characteristic: the resonance structure, the spin of the resonance and the flavor composition of the decays[19], can help in separating signal from background. In particular, when the Higgs mass enters the Z and W resonance region, flavor tagging is indispensable leading to a sample of about 240 events composed of 60 $W \to jj$, 80 $Z \to jj$ and 100 $H \to jj$ tagged as $b\bar{b}$–jets (again for a luminosity of 20 fb^{-1} and for realistic tagging efficiencies). For a Higgs mass around $2M_W$, the background process with W^+W^- and WZ final states can be reduced to a negligible level.

2.4.3. Radiation off the Top

The coupling of the Higgs boson to other particles is proportional to their masses. Therefore for a heavy top quark the Higgs radiation off the top quark-antiquark pair produced in e^+e^- collisions becomes an interesting process to look at. In general, the top and Higgs masses must be kept non–zero in the calculations so that the cross section for Higgs bremsstrahlung is quite involved. It can be found in Ref. [17,21] The integrated cross section is shown for various c.m. energy values in Fig. 5 as a function of M_H. At $\sqrt{s} = 500$ GeV, while for small M_H the cross sections increase with m_t as a result of the rising Yukawa coupling, this trend is reversed for heavy Higgses by

Fig. 5 The cross sections $\sigma(e^+e^- \to t\bar{t}H)$ at $\sqrt{s} = 0.5, 1$ and 1.5 TeV as a function of the Higgs mass; the top quark mass is fixed to 175 GeV.

the reduction of the available phase space. For an integrated luminosity of $\int \mathcal{L} = 20$ fb^{-1}, some 100 events can be expected at Higgs masses of order 60 GeV, falling to less than 20 events at 100 GeV. With such a small sample of events the process $e^+e^- \to t\bar{t}H$ is not suitable for Higgs search. However an interesting point is that the $t\bar{t}H$ final state is generated almost *exclusively* through Higgs bremsstrahlung off the top quarks. There is an additional contribution with Higgs particles emitted by the Z line followed by $Z \to t\bar{t}$ which turns out to be small, of the order $1-2$ %. Therefore this process may allow for a *direct* measurement of Yukawa coupling of the Higgs to top quarks. There is a reasonable hope to isolate these events experimentally despite the low rates since the signature of the process $e^+e^- \to t\bar{t}H \to WWbbbb$ is spectacular. The large number of b quarks together with the mass constraints $m(bb) = m(H)$ and $m(Wb) = m(t)$ will be crucial in rejecting background events.

Even though Higgs bremsstrahlung off top quarks is not easy to handle experimentally in view of small cross sections, it nevertheless deserves attention as it may provide the opportunity to measure directly the Higgs–fermion coupling.

2.5. Higgs production at future $\gamma\gamma$ colliders

Since the photons couple to Higgs bosons via heavy particle loops, the $H\gamma\gamma$ amplitudes are sensitive to all charged particles, standard and nonstandard, well above the Higgs masses themselves. The cross section for the $\gamma\gamma$ fusion of Higgs bosons is found by folding the parton cross section with the $\gamma\gamma$ luminosity, see e.g. Ref.[22,23] and the parton cross section is determined by the $\gamma\gamma$ width ($s_{\gamma\gamma}$ is the photon c.m. energy squared),

$$\sigma(\gamma\gamma \to \Phi) = \frac{8\pi^2}{M_\Phi^3}\Gamma(\Phi \to \gamma\gamma)\delta(1 - M_\Phi^2/s_{\gamma\gamma}) \tag{12}$$

With increasing energy of e^+e^- collider the $\gamma\gamma$ luminosity of bremsstrahlung photons increases and $\gamma\gamma$ fusion may become an important production process. For example, assuming integrated luminosity of 10 fb^{-1} at $\sqrt{s} = 2$ TeV e^+e^- collider (CLIC) it was found[23] that some 20 - 30 events can be expected for M_H in the range 50 - 250 GeV, falling quickly to 1 at 400 GeV. Here relatively low rates are due to the

fact that for bremsstrahlung photons the luminosity is still low and falls very quickly with the energy available in $\gamma\gamma$ collisions.

The main interest in $\gamma\gamma$ fusion, however, stems from the fact that future high–energy e^+e^- linear colliders can be *converted* to a high energy $\gamma\gamma$ colliders,[24] by using Compton back–scattered laser light. In this way one converts electrons and positrons to energetic photons with *practically* the same total energy ($\sim 80\%$) and luminosity ($\sim 100\%$) as the the original e^+e^- collider.

Studies performed in Ref.[25] show that in a $\gamma\gamma$ collider one can expect as many as 10^4 events per year. Unfortunalely there will be a copious production of $b\bar{b}$ and WW pairs $\sim 10^5 - 10^6$/year as well which will complicate Higgs search. However once the Higgs boson is found and its mass measured, detailed studies of its properties can be undertaken by tuning the energy of $\gamma\gamma$ collider to M_H. Operating with photons of given polarization may help to suppress the background.

In particular, we will see in section 3 that in the e^+e^- mode, the lightest supersymmetric Higgs particle h can be detected but it cannot be distinguished from the $\mathcal{SM}$ Higgs boson in some part of the $\mathcal{MSSM}$ parameter space if $SUSY$ decays are not allowed. One of the ways to distinguish h from the $\mathcal{SM}$ Higgs particle can be provided by Higgs production in $\gamma\gamma$ fusion. In the Standard Model this process is built up by W and top quark loops which interfere destructively. Additional contributions, (provided for example by supersymmetric particles: chargino, sfermion and charged Higgs boson loops) can alter the $\mathcal{SM}$ production rates and thus help in revealing the nature of Higgs boson.

3. Search for the Higgs bosons in supersymmetric extension of the Standard Model

3.1. The Higgs sector of the $\mathcal{MSSM}$

The minimal supersymmetric extension $\mathcal{MSSM}$[1] employs only two doublets Φ_1 and Φ_2. The field Φ_2 couples only to up–type quarks while Φ_1 couples to down–type quarks and charged leptons. The physical Higgs bosons introduced in this extension are of the following type: two $\mathcal{CP}$–even neutral bosons h and H (where by convention $M_h \leq M_H$), a $\mathcal{CP}$–odd neutral boson A (usually called pseudoscalar) and two charged Higgs bosons $H^\pm$.

The properties of the scalar particles and their interactions with gauge bosons and fermions depend on Higgs masses, M_h, M_H, M_A and $M_{H^\pm}$, and two mixing angles: α in the $\mathcal{CP}$–even neutral, and β in the charged Higgs sectors. The angle β is related to the ratio of the vacuum expectation values $\tan\beta = v_2/v_1$ (where v_1 (v_2) is the vaccum expectation value of the field Φ_1 (Φ_2)). However supersymmetry leads to several relations among these parameters and, in fact, only two of them are independent at tree level. These relations impose a strong hierarchical structure on the mass spectrum ($M_h \leq M_Z, M_A \leq M_H$ and $M_W \leq M_{H^\pm}$) which however is broken by radiative corrections[26] since the top quark mass is large. The parameter $\tan\beta$ will be assumed in the range $1 < \tan\beta < m_t/m_b$, consistent with restrictions that follow from interpreting the $\mathcal{MSSM}$ as a low energy limit of a supergravity model.

Since the lightest $\mathcal{CP}$–even scalar boson h is likely to be discovered first, an attractive choice of the two input parameters is the set $(M_h, \tan\beta)$. Once these two parameters (and the top quark mass and the associated squark masses which enter through radiative corrections) are specified, all other quantities are predicted. The most important radiative corrections[26] can be determined by the parameter ϵ

$$\epsilon = \frac{3\alpha}{2\pi}\frac{1}{s_W^2 c_W^2}\frac{1}{\sin^2\beta}\frac{m_t^4}{M_Z^2}\log\left(1 + \frac{M_S^2}{m_t^2}\right) \tag{13}$$

where $s_W^2 = 1 - c_W^2 \equiv \sin^2\theta_W$. These corrections shift the mass of the light neutral Higgs boson h upward with increasin top mass. The upper limit on M_h as a function of top quark mass is shown in Fig. 6 for squark masses $M_S = 1$ TeV and two representative values of $\tan\beta = 2.5$ and 20; and update of Ref.[27] The upper bound on M_h is shifted from the tree level value M_Z up to ~ 130 GeV for $m_t = 175$ GeV and ~ 140 GeV for $m_t = 200$ GeV.

Fig. 6 The masses of the Higgs particles in the $\mathcal{MSSM}$ including radiative corrections. a) Upper limit on M_h as a function of m_t; b)–d) masses of the H, A and $H^\pm$ Higgs bosons as functions of M_h with the top mass fixed to 175 GeV. The dashed curve shows the leading correction, $A_t = A_b = \mu = 0$, while the solid curves include the full corrections, $A_t = A_b = 1$ TeV and $\mu = -200, 0, 200$ GeV; μ is a $SUSY$ Higgs mass and A_t, A_b are soft $SUSY$ breaking couplings.

It is interesting to note that for $m_t = 175$ GeV the upper limit for M_h (≤ 135 GeV) is smaller than the lower limit of $\mathcal{SM}$ Higgs mass (≥ 160 GeV) – see Fig. 1. Therefore the measurement of M_h might give the first indication which model is realized in Nature.

Taking M_h and $\tan\beta$ as the input parameters, masses of other Higgses are given by

$$
\begin{aligned}
M_A^2 &= \frac{M_h^2(M_Z^2 - M_h^2 + \epsilon) - \epsilon M_Z^2 \cos^2\beta}{M_Z^2 \cos^2 2\beta - M_h^2 + \epsilon \sin^2\beta} \\
M_H^2 &= M_A^2 + M_Z^2 - M_h^2 + \epsilon \\
M_{H^\pm}^2 &= M_A^2 + M_W^2
\end{aligned}
\tag{14}
$$

In the subsequent discussion, we will use for definiteness the values $m_t = 175$ GeV and $M_S = 1$ TeV. For the two representative values of $\tan\beta$ introduced above, the masses M_A, M_H and $M_{H^\pm}$ are displayed in Fig. 9b–d as a function of the light neutral Higgs mass M_h. Notice that for $M_h < M_h^{max}$ for a given value of $\tan\beta$, the Higgs masses are of the order 100 to 200 GeV for M_H and $M_{H^\pm}$, and up to ~ 150 GeV for M_A (similarly to M_h). On general grounds, the masses of the heavy neutral and charged Higgs bosons are expected to be of the order of the electroweak symmetry breaking scale.

The mixing angle α is expressed as

$$
\mathrm{tg}2\alpha = \mathrm{tg}2\beta \frac{M_A^2 + M_Z^2}{M_A^2 - M_Z^2 + \epsilon/\cos 2\beta} \qquad \left[-\frac{\pi}{2} \leq \alpha \leq 0 \right]
\tag{15}
$$

The couplings of Higgs bosons to fermions and gauge bosons in general depend on the angles α and β. Normalized to the $\mathcal{SM}$ Higgs couplings, they are summarized in Table 1.

Tab. 1: Higgs bosons couplings in the $\mathcal{MSSM}$ to up-type and down-type fermions and gauge bosons (V=W,Z) relative to the $\mathcal{SM}$ Higgs boson couplings.

Φ	$g_{\Phi\bar{u}u}$	$g_{\Phi\bar{d}d}$	$g_{\Phi VV}$
H_{SM}	1	1	1
h	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$	$\sin(\beta - \alpha)$
H	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$	$\cos(\beta - \alpha)$
A	$1/\tan\beta$	$\tan\beta$	0

For $\tan\beta > 1$ the couplings to down (up) type fermions are enhanced (suppressed) compared to the $\mathcal{SM}$ Higgs couplings. If M_h is very close to its upper limit for a given value of $\tan\beta$, the couplings of h to fermions and gauge bosons are $\mathcal{SM}$ like. It may therefore be very difficult to distinguish the Higgs sector of the $\mathcal{MSSM}$ from the $\mathcal{SM}$, if all other Higgs bosons are very heavy[28].

Comment on LEP 1 results for $\mathcal{MSSM}$ and 2$\mathcal{HDM}$ Higgs bosons

As we already noticed, in $\mathcal{MSSM}$ there are only two independent parameters in the Higgs sector, for example M_h and $\tan\beta$. At LEP 1 energies the heavier $\mathcal{CP}$ Higgs

boson cannot be produced on-shell ($M_H > M_Z$) and there are only two important production mechanisms for h and/or A: $Z \to Z^*h$ and $Z \to Ah$. Negative results of $SUSY$ Higgs particles searches at LEP in these processes exclude[29] h and A bosons with masses smaller than $M_{h,A} \simeq 45$ GeV for $m_t = 140$ GeV, $M_S = 1$ TeV and tg$\beta > 1$. If the parameters of the model are allowed to vary arbitrarily and if one includes all possible decay modes, this bound becomes[29] $M_h > 44$ GeV and $M_A > 21$ GeV.

In a general two-Higgs doublet model the Higgs mass regions excluded experimentally are much more model dependent. The reason is that without supersymmetry the masses of Higgs bosons (M_h, M_H, M_A and $M_{H^\pm}$) and mixing angles (α and β) are arbitrary and unrelated parameters. Therefore the other $\mathcal{CP}$-even Higgs boson H can be produced if light enough. In addition other production processes, namely Higgs (h, H or A) radiation off the b-quark or τ-lepton pairs from Z decays can be important[30]. For example for large tanβ the couplings of h and A are enhanced to down-type fermions. Therefore if ZZh is supressed by mixing angles and the Higgs pair Ah production is supressed kinematically then the radiation of h or A (whichever is lighter) off b-quarks or τ-leptons becomes dominant[31]. (In $SUSY$ model for large tanβ and h light enough to be produced the CP-odd A is almost mass degenerate with h so Higgs pair hA production process is kinematically possible and dominates.) This is illustrated in Fig. 7, where contour lines for cross sections of processes (a) $e^+e^- \to hZ \to hb\bar{b}$ and (b) $e^+e^- \to b\bar{b} \to hb\bar{b}$ are shown as functions of M_h and tanβ and for a specific choice of other parameters. Experimentally the radiation processes, which have a different signature, have not been looked for. Before these processes are analysed we cannot conclude that LEP 1 excludes $2\mathcal{HDM}$ Higgs bosons within kinematical reach in a model independent way.

Fig. 7 Regions in the plane (M_h, tanβ=v_2/v_1) corresponding to the dominant process (a) or (b). Contour lines are drawn for $\log_{10}(\sigma/1\text{pb})$. Solid (dashed) lines are for $\epsilon = 0.2(2)$. For details, see Ref.[31]

3.2. Decay modes

3.2.1. Decays to $\mathcal{SM}$ particles

The partial decay width of a neutral Higgs boson Φ into fermion pairs is given by

$$\Gamma(\Phi \to \bar{f}f) \;=\; N_c \frac{G_F m_f^2}{4\sqrt{2}\pi}\, g_{\Phi ff}^2\, M_\Phi\, \beta^p \tag{16}$$

where $\beta = (1 - 4m_f^2/M_\Phi^2)^{1/2}$ and $p = 3(1)$ for the $\mathcal{CP}$–even (odd) Higgs boson; the couplings $g_{\Phi ff}$ are listed in Tab.1. For final state quarks one has to include QCD corrections. As in $\mathcal{SM}$ the use of the running masses takes into account the bulk of these corrections, which in the limit $M_H \gg m_q$ are the same for $\mathcal{CP}$–odd and $\mathcal{CP}$–even Higgs bosons.

The lightest Higgs boson will decay mainly into fermion pairs since its mass is smaller than ~ 130 GeV. This is also the dominant decay mode of the pseudoscalar boson A which has no tree level couplings to gauge bosons. For values of $\tan\beta$ larger than unity and for masses less than ~ 130 GeV, the main decay modes of h and A will be decays into $b\bar{b}$ and $\tau^+\tau^-$ pairs with the branching ratios being always larger than $\sim 90\%$ and 5%, respectively. The decays into $c\bar{c}$ and gluons are in general strongly suppressed especially for large values of $\tan\beta$. For large masses, the top decay channels $H, A \to t\bar{t}$ open up; yet this mode remains suppressed for large $\tan\beta$.

If the mass is high enough, the heavy $\mathcal{CP}$–even Higgs boson can in principle decay into weak gauge bosons $H \to VV$, $V = W$ or Z. Below the threshold it can decay into VV^* pairs with one of the vector bosons being virtual. (For h this decay mode is negligible.) The partial decay width is given by

$$\Gamma(H \to VV^{(*)}) = g_{HVV}^2 \Gamma(H_{\rm SM} \to VV^{(*)}) \tag{17}$$

where $\Gamma(H_{\rm SM} \to VV^{(*)})$ given by Eqs. (8,9,10). Since the H partial width is proportional to $\cos^2(\beta-\alpha)$, it is strongly suppressed because if M_H is large enough for these decay modes to be kinematically allowed, M_h is very close to its maximum so that $\cos^2(\beta-\alpha) \to 0$. For the same reason, the decay $A \to Zh$, is also suppressed[1]

$$\Gamma(A \to Zh) = \frac{G_F}{8\sqrt{2}\pi}\cos^2(\beta-\alpha)\,\frac{M_Z^4}{M_A}\lambda^{1/2}(M_Z^2, M_h^2; M_A^2)\lambda(M_A^2, M_h^2; M_Z^2) \tag{18}$$

with $\lambda(x, y; z) = (1 - x/z - y/z)^2 - 4xy/z^2$ being the usual two–body phase space function.

The heavy neutral Higgs boson H can also decay into two lighter Higgs bosons h or A,[1]

$$\Gamma(H \to hh) \;=\; \frac{G_F}{16\sqrt{2}\pi}\frac{M_Z^4}{M_H}\left(1 - 4\frac{M_h^2}{M_H^2}\right)^{1/2}[\cos 2\alpha\cos(\beta+\alpha) - 2\sin 2\alpha\sin(\beta+\alpha)]^2$$

$$\Gamma(H \to AA) \;=\; \frac{G_F}{16\sqrt{2}\pi}\frac{M_Z^4}{M_H}\left(1 - 4\frac{M_A^2}{M_H^2}\right)^{1/2}[\cos 2\beta\cos(\beta+\alpha)]^2 \tag{19}$$

These modes, however, are restricted to small domains in the parameter space.

Decays $\Phi \to gg$ are mediated by top and bottom quark loops since the squarks decouple from the effective Φgg vertex for high masses. This decay mode is small. The branching ratio reaches a few % for h-decay only for M_h close to the maximum value where h has $\mathcal{SM}$ like couplings, and for H-decay only below 140 GeV and small values of $\tan\beta$ where the coupling to top quarks is sufficiently large. For the pseudoscalar Higgs particle, the gluonic decay mode is also small. Decays of Higgs bosons to $\gamma\gamma$ and $Z\gamma$ final states are very rare with branching ratios of order $\mathcal{O}(10^{-3})$ or below.

The coupling of the charged Higgs particle to fermions is a mixture of scalar and pseudoscalar couplings

$$g_{H^+ u\bar{d}} = \left(\frac{G_F}{\sqrt{2}}\right)^{1/2} \left[(1 - \gamma_5)\frac{m_u}{\tan\beta} + (1 + \gamma_5)m_d\tan\beta\right] \tag{20}$$

The charged Higgs particles decay into fermions with a partial decay width

$$\Gamma(H^+ \to u\bar{d}) = \frac{N_c G_F \lambda^{\frac{1}{2}}}{4\sqrt{2}\pi M_{H^\pm}} \left[(M_{H^\pm}^2 - m_u^2 - m_d^2)\left(m_d^2 \mathrm{tg}^2\beta + \frac{m_u^2}{\mathrm{tg}^2\beta}\right) - 4m_u^2 m_d^2\right] \tag{21}$$

with $\lambda^{\frac{1}{2}} = \lambda^{\frac{1}{2}}(m_u^2, m_d^2; M_{H^\pm}^2)$ and, if allowed kinematically, they also decay into the lightest neutral Higgs plus a W boson,

$$\Gamma(H^+ \to Wh) = \frac{G_F \cos^2(\beta - \alpha)}{8\sqrt{2}\pi c_W^2} \frac{M_W^4}{M_{H^\pm}} \lambda^{\frac{1}{2}}(M_W^2, M_H^2; M_{H^\pm}^2)\lambda(M_{H^\pm}^2, M_h^2; M_W^2) \tag{22}$$

Fig. 8 Total decay widths of the $SUSY$ Higgs bosons (without decays into $SUSY$ particles) as functions of their masses for (a) $\mathrm{tg}\beta = 2.5$ and (b) $\mathrm{tg}\beta = 20$. The top mass was chosen as $m_t = 175$ GeV and $M_S = 1$ TeV.

Below the tb and Wh thresholds, the charged Higgs particles will decay mostly into $\tau\nu_\tau$ and $c\bar{s}$ pairs, the former being dominant for $\tan\beta > 1$. For large $M_{H^\pm}$ values, the top–bottom decay $H^+ \to t\bar{b}$ becomes dominant.

Adding up the various decay modes,[27,32] the width of all five Higgs bosons remains very small, even for large masses. This is shown for the two representative values $\tan\beta = 2.5$ and 20 in Fig. 8. Apart from the $\mathcal{CP}$–even heavy neutral Higgs boson H and small $\tan\beta$, the pattern of branching ratios is in general quite simple. The neutral Higgs bosons decay preferentially to $b\bar{b}$, and to a lesser extent to $\tau^+\tau^-$ pairs; the charged Higgs bosons to $\tau\nu_\tau$ and, preferentially, $t\bar{b}$ pairs above this threshold.

3.2.2. Decays to supersymmetric particles

At least for the heavy Higgs bosons H, A and $H^\pm$ decays into charginos and neutralinos could eventually play a significant role since some of these particles are expected to be light enough. From the negative search of supersymmetric particles in Z decays,[36] the lightest neutralino $[\tilde{\chi}_1^0]$ mass is restricted to be larger than 20 GeV for $\tan\beta = 2.5$ and larger than 22 GeV for $\tan\beta > 4$; the second lightest neutralino $[\tilde{\chi}_2^0]$ and the charginos are excluded if their masses are less than $\sim M_Z/2$. If the search at LEP200 with a c.m. energy of 180 GeV is negative, charginos with masses $m_{\tilde{\chi}_1^+} < 90$ GeV will also be excluded.

Fig. 9. Contour lines in the (μ, M) plane where the sum of the branching ratios of the lightest Higgs boson h into charginos and neutralinos exceeds 5% (dashed lines) and 50% (full lines) for two values of M_h and $\tan\beta$; the shaded areas are the regions which are (can be) excluded at LEP 1 (LEP200). The top mass was taken to be $m_t = 140$ GeV.

Sfermions are probably too heavy to affect Higgs decays. The present LEP data exclude sleptons with masses below $\sim M_Z/2$; if sleptons will not be observed at LEP200 these limits can be improved by roughly a factor of two. On the other hand, CDF data[33] restrict the squarks masses to be larger than ~ 150 GeV if cascade decays are suppressed. We shall assume in this discussion that squarks and sleptons are heavy so that they will not affect Higgs boson decays and production.

The decay widths of the neutral and charged Higgs bosons into chargino or neutralino pairs can be found in Ref.[34] In Fig. 9 (an update of Ref.[27,35]) the contour lines are shown in the (μ, M) plane where the sum of the branching ratios of h into the lightest chargino pair and the lightest and next–to–lightest neutralino pairs exceeds 5% (dashed lines) and 50% (solid lines). (M is a universal gaugino mass.) These decays can be important for h masses close to the maximum allowed values; in this case the lightest Higgs boson has $\mathcal{SM}$ couplings and the dominant $b\bar{b}$ decay mode is not enhanced anymore for large $\tan\beta$ values, so that other decay modes can become significant. For these masses and for large $\tan\beta$ values, the branching ratios for neutralino/chargino decays are sizable even outside the regions which can be probed at LEP200.

The decays of H, A or $H^{\pm}$ into chargino and neutralino pairs can be very large.[27] They can exeed 50% in some areas of the $\mathcal{MSSM}$ parameter space, $i.e.$ for positive values of μ and/or M values below 200 GeV. This is due to the fact that the couplings of the Higgs bosons to charginos and neutralinos are gauge couplings which can be larger than the couplings to standard fermions and gauge bosons. Finally, the branching fractions of the invisible neutral Higgs decays can be important, [14,27] and they could jeopardize the search for the Higgs particles at hadron colliders. However, at e^+e^- colliders missing mass techniques can be used to detect invisibly decaying Higgs particle produced in association with the Z boson or in mixed visible and invisible decays modes in Ah or HA production.

3.3. Production mechanisms of $\mathcal{MSSM}$ Higgs bosons at future e^+e^- colliders

3.3.1. Neutral Higgs Bosons
The main production mechanisms of neutral Higgs bosons at e^+e^- colliders are the bremsstrahlung, pair production and fusion processes

$$i) \quad \text{bremsstrahlung} \quad e^+e^- \;\to\; (Z) \to Z + h/H \tag{23}$$

$$ii) \quad \text{pair production} \quad e^+e^- \;\to\; (Z) \to A + h/H \tag{24}$$

$$iii) \quad \text{fusion processes} \quad e^+e^- \;\to\; \nu\,\bar{\nu}\,(WW) \to \nu\,\bar{\nu}\; + h/H$$

$$e^+e^- \;\to\; e^+e^-(ZZ) \to e^+e^- + h/H \tag{25}$$

The $\mathcal{CP}$–odd Higgs boson A cannot be produced in fusion processes to leading order. The Higgs radiation off top quark is supressed for $\tan\beta > 1$.

The cross sections for the production processes i) and ii) can be written as

$$
\begin{aligned}
\sigma(e^+e^- \to Zh) &= \sin^2(\beta - \alpha)\,\sigma_{\text{SM}} \\
\sigma(e^+e^- \to ZH) &= \cos^2(\beta - \alpha)\,\sigma_{\text{SM}} \\
\sigma(e^+e^- \to Ah) &= \cos^2(\beta - \alpha)\,\sigma_{\text{SM}}\,\bar{\lambda} \\
\sigma(e^+e^- \to AH) &= \sin^2(\beta - \alpha)\,\sigma_{\text{SM}}\,\bar{\lambda}
\end{aligned}
\tag{26}
$$

where σ_{SM} denotes the cross section for Higgs bremsstrahlung in $\mathcal{SM}$, Eq. (10), and

$$
\bar{\lambda} = \frac{\lambda^{3/2}(M_j^2, M_A^2; s)}{\lambda^{1/2}(M_j^2, M_Z^2; s)\,(\lambda(M_j^2, M_A^2; s) + 12 M_Z^2/s)} \qquad (j = h, H) \tag{27}
$$

Fig. 10 Production cross sections of the $\mathcal{CP}$–even neutral Higgs bosons at $\sqrt{s} = 500$ GeV as functions of their masses for three values of $\mathrm{tg}\beta = 2.5, 5$ and 20; (a) Bremsstrahl processes $e^+e^- \to Z + h/H$, and (b) in association with the pseudoscalar Higgs boson $e^+e^- \to A + h/H$.

accounts for the correct suppression of the P–wave cross sections near the threshold. Since the cross sections in Eq. (26) come with either with a coefficient $\sin^2(\beta - \alpha)$ or $\cos^2(\beta - \alpha)$ they are mutually complementary to each other. Therefore at least one of the $\mathcal{CP}$–even Higgs bosons should be detected because σ_{SM} is large. Let us discuss them in more detail.

The cross sections for the production of the bosons h and H via bremsstrahlung are shown as functions of the Higgs mass in Fig. 10a. The cross section for h is large for small values of $\tan\beta$ and/or large values of M_h where $\sin^2(\beta - \alpha)$ approaches maximum value. In both cases the cross section is of the order of ~ 50 fb, which for an integrated luminosity of 20 fb^{-1} corresponds to ~ 1000 events. On the other hand, the cross section for H is large for large $\tan\beta$ and light h (implying small M_H). In the case of h (and also for H in most of the parameter space) the signal consists of a Z boson accompanied by a $b\bar{b}$ or a $\tau^+\tau^-$ pair. The signal is easy to separate from the background which comes mainly from ZZ production if the Higgs mass is close to M_Z. Results for the associate production $e^+e^- \to Ah$ and AH are displayed in Fig. 10b. The situation is opposite to the previous case: for light h and/or large values of $\tan\beta$ the cross section for Ah is large whereas AH production is preferred in the complementary region. The sum of the two cross sections decreases from ~ 50 to 10 fb if M_A increases from ~ 50 to 200 GeV. Signals mainly consist of four b quarks in the final state, requiring efficient b quark tagging. Mass constraints will help to eliminate the backgrounds from QCD jets as well as ZZ final states. The above discussion is summarized in Fig. 11, from Ref.[14], an update of Ref.[27]

Fig. 11 Regions of the $(M_h, \tan\beta)$ plane where $e^+e^- \to hZ, hA, HZ$ and HA are observable, *i.e.* the cross sections are larger than 2.5 fb. The dashed area is theoretically forbidden. The region to the left of the thin line can be probed at LEP200. The process $e^+e^- \to hZ$ is accessible in the entire area below the full line, hA in the entire area above the broken line and HZ in the entire area above the full line; HA final states can be detected in the area between the two dashed lines.

WW and ZZ fusion provide additional mechanisms for the production of the $\mathcal{CP}$–even neutral Higgs bosons. They can again be expressed in terms of the corresponding $\mathcal{SM}$ cross sections, given in Eq. (11)

$$
\begin{aligned}
\sigma(e^+e^- \to (VV) \to h) &= \sin^2(\beta - \alpha)\sigma_{\mathrm{SM}}^{VV} \\
\sigma(e^+e^- \to (VV) \to H) &= \cos^2(\beta - \alpha)\sigma_{\mathrm{SM}}^{VV}
\end{aligned}
\tag{28}
$$

As in the case of the bremsstrahlung process, the production of light h and heavy H bosons are complementary.

Fig. 12 Production cross sections of the $\mathcal{CP}$–even neutral Higgs bosons at $\sqrt{s} = 500$ GeV as functions of their masses for three values of $\mathrm{tg}\beta = 2.5, 5$ and 20; (a) WW fusion process $e^+e^- \to \nu\bar{\nu} + h/H$ (b) and ZZ fusion process $e^+e^- \to e^+e^- + h/H$.

For the Higgs mass less than 160 GeV at $\sqrt{s} = 500$ GeV the WW fusion mechanism has cross sections larger than the bremsstrahlung process. However, the final state cannot be fully reconstructed and the signal is more difficult to extract. The cross sections for the ZZ fusion mechanism are about an order of magnitude smaller than for the one for WW. Nevertheless it will be useful as the final state can be fully reconstructed. For representative values of $\tan\beta$ they are shown in Fig. 12.

3.3.2. Charged Higgs Bosons

An unambiguous signal of an extended Higgs sector would be the discovery of a charged Higgs boson. In a general two–Higgs doublet model, charged Higgs bosons can be as light as ~ 45 GeV, the lower limit derived from the negative search at LEP 1.[36] In the $\mathcal{MSSM}$ however, $H^\pm$ is constrained to be heavier than the W boson. More precisely, the lower limit $M_A > 45$ GeV obtained at LEP 1 implies $M_{H^\pm} > 90$ GeV.

The production of a pair of charged Higgs bosons in e^+e^- collisions proceeds through virtual photon and Z boson exchange.

$$\sigma(e^+e^- \longrightarrow H^+H^-) = \frac{\pi\alpha^2}{3s} \left[1 - \frac{2\hat{v}_e\hat{v}_H}{1 - M_Z^2/s} + \frac{(\hat{a}_e^2 + \hat{v}_e^2)\hat{v}_H^2}{(1 - M_Z^2/s)^2} \right] \beta^3 \tag{29}$$

with the standard Z charges $\hat{v}_e = (-1 + 4s_W^2)/4c_W s_W$, $\hat{a}_e = -1/4c_W s_W$ and $\hat{v}_H = (-1 + 2s_W^2)/2c_W s_W$, and $\beta = (1 - 4M_{H^\pm}^2/s)^{1/2}$. The cross section is shown in Fig. 13a as a function of the charged Higgs mass for a c.m. energy $\sqrt{s} = 500$ GeV (it does not depend on any extra parameter). For small Higgs masses the cross section is of order 100 fb, but it drops very quickly due to the P–wave suppression factor β^3 near the threshold. For $M_{H^\pm} = 220$ GeV, the cross section has fallen to a level of $\simeq 5$ fb, which for an integrated luminosity of 20 fb^{-1} corresponds to 100 events. The angular distribution of the charged Higgs bosons follows the $\sin^2\theta$ law typical for spin–zero particle production.

Fig. 13 Production cross sections of the charged Higgs boson (a) at e^+e^- colliders with $\sqrt{s} = 500$ GeV and in $\gamma\gamma$ collisions with $\sqrt{s} = 400$ and (b) in top decays with $m_t = 175$ GeV and for three values of $\tan\beta = 2.5, 5$ and 20.

The charged Higgs boson, if lighter than the top quark, can also be produced in top decays.[37] In the range $1 < \tan\beta < m_t/m_b$ favored by $\mathcal{SUSY}$ models, the branching ratio BR($t \to bH^+$) varies between $\sim 2\%$ and 20%. Since the cross section for top pair production is of order of 50 fb at $\sqrt{s} = 500$ GeV, this corresponds to 20 and 200 charged Higgs bosons at a luminosity $\int \mathcal{L} = 20$ fb^{-1}; Fig. 13b.

If $M_{H^\pm} < m_t + m_b$, the charged Higgs boson will decay mainly into $\tau\nu_\tau$ and $c\bar{s}$ pairs, the $\tau\nu_\tau$ mode dominating for $\tan\beta$ larger than unity. This results in a surplus of τ final states over e, μ final states, an apparent breaking of τ $vs.$ e, μ universality. For large Higgs masses the dominant decay mode is the top decay $H^+ \to t\bar{b}$. In some part of the parameter space also the decay $H^+ \to W^+h$ is allowed, leading to cascades with heavy τ and b particles in the final state.

3.4. $\mathcal{MSSM}$ Higgs bosons in $\gamma\gamma$ Collider

The scalar $\mathcal{MSSM}$ Higgs particles h, H couple to vector bosons directly, and therefore the positive parity can be checked by analyzing the Z final states in $e^+e^- \to Z^* \to ZH$ ($Z \to f\bar{f}$) as discussed above. However, the pseudoscalar A boson has no tree level couplings to the vector bosons and the latter methods cannot be used. To study[38] the $\mathcal{CP}$ properties of both scalar and pseudoscalar Higgs bosons on equal footing, one can run in the $\gamma\gamma$ mode where both type of particles are produced through loop diagrams with similar rates. For scalar particles the production amplitude is maximal only for parallel vectors while pseudoscalars prefer perpendicular polarization vectors. It remains to be seen whether the degree of linear polarization of the beams will be high enough to be efficiently explored.

Charged Higgs particles can also be created in $\gamma\gamma$ collisions. The cross section is given by

$$\sigma(\gamma\gamma \to H^+H^-) = \frac{2\pi\alpha^2}{s}\beta\left[2 - \beta^2 - \frac{1-\beta^4}{2\beta}\log\frac{1+\beta}{1-\beta}\right] \qquad (30)$$

where β is the velocity of the Higgs particle. The numerical result is displayed in Fig. 13a for the $\gamma\gamma$ luminosity without beam polarization. Due to the reduced energy, the maximum Higgs mass which can be probed in $\gamma\gamma$ collisions is smaller than in the original e^+e^- collisions; the cross section however is enhanced by a factor ~ 3 in the low mass range.

4. Summary

The Higgs search and tests of the electroweak symmetry breaking will be the most important physics goals of future high–energy colliders. In this review, we have discussed the properties of the Higgs particles of the Standard Model and of its minimal supersymmetric extension. We have updated the new value of the top quark various results for Higgs couplings, decay widths and production cross sections in future e^+e^- and $\gamma\gamma$ colliders.

e^+e^- linear colliders with energies in the range $300 - 500$ GeV and a luminosity of a few times 10^{33}cm^{-2}s^{-1} are ideal machines to search for Higgs bosons in the mass range below the scale of electroweak symmetry breaking.

In the intermediate mass range, Standard Higgs particles can be observed in three independent production channels: the bremsstrahlung process $e^+e^- \rightarrow ZH$ and the fusion processes $e^+e^- \rightarrow \bar{\nu}\nu H$ and $e^+e^- \rightarrow e^+e^- H$. The particle is relatively easy to detect especially in the ZH channel, where the main background from ZZ pair production can be suppressed efficiently by using micro–vertex detectors since Higgs bosons with masses below 140 GeV decay mainly into $b\bar{b}$ pairs.

An even stronger case for e^+e^- colliders operating in the 500 GeV range is made by supersymmetric extensions of the Standard Model. Since in $\mathcal{MSSM}$ the lightest Higgs particle has a mass below 140 GeV and decays mainly into $b\bar{b}$ and $\tau^+\tau^-$ pairs, it cannot be missed at an e^+e^- collider with an energy $\sqrt{s} > 300$ GeV, independently of its decay modes and in the entire $SUSY$ parameter space. The heavy neutral Higgs particles can be produced in the bremsstrahlung and fusion processes or pairwise, these processes being complementary. At least one neutral Higgs boson must be detected or $\mathcal{MSSM}$ rejected. In a large part of the $SUSY$ parameter space, all neutral Higgs bosons can be observed. Charged Higgs particles can be detected up to practically the kinematical limit.

Once the Higgs boson is found, its fundamental properties can be investigated. The Higgs spin can be measured by analyzing the angular dependence of the ZH production process and in the Higgs decays into massive gauge bosons. The Higgs couplings to the massive gauge bosons can be determined through the production rates, the coupling to heavy fermions through the Higgs decay branching fractions, and in some mass window, Higgs radiation off top quarks.

Future e^+e^- can be turned to very high–energy $\gamma\gamma$ or $e\gamma$ colliders by using back–scattering of laser light. The $\gamma\gamma$ mode of the e^+e^- collider could be useful to measure accurately the Higgs–photons coupling to which new particles might contribute, and to study the $\mathcal{CP}$ properties of the Higgs particles.

Although we did not discuss the hadron colliders in this review, it turns out that e^+e^- linear colliders operating in the 300–500 GeV energy range and hadron colliders operating in the multi–TeV range have a complementary potential[7] for exploring the key issue of the mechanism of electroweak symmetry breaking.

Acknowledgments

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